Electroweak Symmetry Breaking without the μ^2 Term

Florian Goertz

Theory Division,
CERN, 1211 Geneva 23, Switzerland

E-mail: florian.goertz@cern.ch

ABSTRACT: We demonstrate that from a low energy perspective a viable electroweak symmetry breaking can be achieved without the (negative sign) μ^2 mass term in the Higgs potential, thereby avoiding completely the appearance of relevant operators. We show that such a setup is self consistent and not ruled out by Higgs physics. In particular, we point out that it is the lightness of the Higgs boson that allows for the electroweak symmetry to be broken dynamically via operators of $D \geq 4$, consistent with the power expansion. Beyond that, we entertain how this scenario might even be preferred phenomenologically compared to the ordinary mechanism of electroweak symmetry breaking, as realized in the Standard Model, and argue that it can be fully tested at the LHC. In an appendix, we classify UV completions that could lead to such a setup, considering also the option of generating all scales dynamically.

Contents

1 Introduction

A common lore, in particular after the discovery of the Higgs boson, is that electroweak symmetry breaking (EWSB) is triggered by a negative sign mass term for a Higgs doublet, either introduced by hand or generated dynamically. The ultraviolet (UV) completion of the Standard Model (SM) is expected to generate (approximately) a Higgs potential of such a form. However, as we will entertain in this article, a more fundamental theory of nature could also have an opposite low energy limit, where the appearance of relevant operators is fully avoided. As we will show, in such a scenario an 'irrelevant' $D = 6$ operator of the type $\mathcal{O}_6 = |H|^6$ could induce a non-trivial vacuum for the scalar sector.[1] We will point out that it is the lightness of the Higgs boson that allows to consider this special setup, where the operators in the Higgs potential are not only deformed in a sub-leading way, as a phenomenologically viable alternative to the SM-form of the potential. It will turn out that this lightness, $m_h^2 \ll v^2$, makes it consistent with the effective field theory (EFT) expansion to fully trade the $(D = 2)$ μ^2-operator for $\mathcal{O}_6$.

After having demonstrated the self-consistency of the setup, we will turn to the phenomenology of the model. The most important collider observable is Higgs-boson pair production, where we will show that the LHC is capable of fully testing

---

[1]Such operators are expected to be present on general grounds since the SM is not anticipated to be the final theory of nature.

the pure version of the idea. Beyond that, the model has intriguing consequences for cosmology. We will see that it just lies in the correct ballpark such as to allow for a strong first order phase transition, as required by electroweak baryogenesis. While in the main part of the paper we just treat the $\mu = 0$ Higgs potential as a distinct and interesting boundary condition that a potential UV model could fulfill, thereby opening up a new direction in model building, in the appendix we will present and classify possible ideas for UV completions of our setup.

2 The Form of the Higgs Potential

We consider the SM without relevant operators and instead augment the Higgs potential with a dimension-6 term $c_6/\Lambda^2 \mathcal{O}_6$ such that it takes the simple form

$$V(H) = \lambda |H|^4 + \frac{c_6}{\Lambda^2} |H|^6 \,, \tag{2.1}$$

where all dimensionfull parameters are either zero or at the cutoff of the theory.

Inspecting the form of the potential, a first observation is that a stable and non-trivial minimum at $|H|^2 > 0$ should be possible if $\lambda < 0$ and $c_6 > 0$. In the following, we will check if such a minimum is also viable phenomenologically. For any given cutoff scale Λ, we can first calculate the position of the minimum, i.e., the vacuum expectation value (vev), denoted as $\langle |H|^2 \rangle \equiv v^2/2$, via $\partial V/\partial |H|^2 = 0$. We find

$$v^2 = -\frac{4}{3} \frac{\lambda}{c_6} \Lambda^2 \,. \tag{2.2}$$

Clearly, the minimization condition only fixes the relative size of the coefficients λ and c_6/Λ^2. In turn, an electroweak-scale vev can be obtained without the need for a large coefficient of the $D = 6$ operator $\mathcal{O}_6$. The size of the latter will however get fixed by the mass of the physical Higgs excitation around the vev, h, where in unitary gauge $H = 1/\sqrt{2}(0, v + h)^T$. This is given as $m_h^2 = \partial^2 V/\partial h^2|_{h=0}$, leading to

$$m_h^2 = 3v^2 \lambda + \frac{15}{4} \frac{c_6}{\Lambda^2} v^4 \,. \tag{2.3}$$

The consequences of these relations will be scrutinized in the next section.

3 Self Consistency of the Setup

We will now examine quantitatively, if it is possible to generate the vev and the Higgs mass at the correct values, keeping the parameters in the range of the validity of the EFT, and then discuss the stability of the low energy setup. To that extent, we first solve eqs. (2.2) and (2.3) for the two free parameters in the potential, λ and
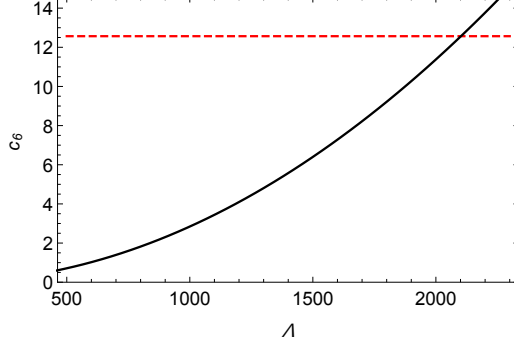
Figure 1. Required value of the Wilson coefficient c_6 in dependence on the cutoff Λ. See text for details.

c_6, expressing them in terms of the vev, fixed by the Fermi constant as $v = 246\,\mathrm{GeV}$, and the Higgs mass $m_h \approx 125\,\mathrm{GeV}$. We obtain

$$\lambda = -\frac{m_h^2}{2v^2} \approx -0.13\,, \quad c_6 = \frac{2m_h^2}{3v^2}\frac{\Lambda^2}{v^2} \approx 2.8\,\frac{\Lambda^2}{\mathrm{TeV}^2}. \tag{3.1}$$

We inspect that, since $m_h^2/v^2 \approx 1/4 \ll 1$, a large cutoff $\Lambda^2 \gg v^2$ is possible while keeping $c_6 \sim \mathcal{O}(1)$. We can thus see explicitly that it is the lightness of the Higgs boson which allows for the mechanism to work naturally. The required c_6 versus the cutoff Λ is visualized in Figure 1, where $c_6 = 4\pi$ is given by the horizontal dashed line for illustration. In particular, setting $\Lambda = 0.8\,\mathrm{TeV}$ ($\Lambda = 1\,\mathrm{TeV}$) requires only $c_6 = 1.8$ ($c_6 = 2.8$) while even $\Lambda = 2\,\mathrm{TeV}$ is still possible in a rather strongly coupled setup with $c_6 = 11.4 < 4\pi$.[2] On the other hand, around $\mathcal{O}(\text{several})$ TeV, at the latest, new physics (NP) would be expected. If the new states are uncolored (which we will assume in the following), such mass scales clearly introduce no tension with current LHC limits. Moreover, we have checked that for all values of the cutoff considered above, the inclusion of a $D = 8$ operator with $\mathcal{O}(1)$ coefficient alters the numerical results by only a few per cent or less. It is thus save to neglect such operators.

We will now study more detailed the correlation between the needed size of the coefficient c_6/Λ^2 and the physical parameters in the Higgs sector, m_h and v, stressing that only a limited part of the larger parameter space, considered before the discovery of the Higgs boson, is viable in our model. To finally remove the redundancy of parameters in the low energy theory, allowing for an easier discussion, we redefine

$$c_6 \equiv \bar{c}_6 \Lambda^2/v^2\,, \tag{3.2}$$

---

[2]Note that UV completions could induce even larger couplings, such as $c_6 \sim 16\pi^2$. Note also that here and in the following we neglect small running effects in λ and c_6 from the cutoff down to the electroweak scale v, where the potential (2.1) is considered.
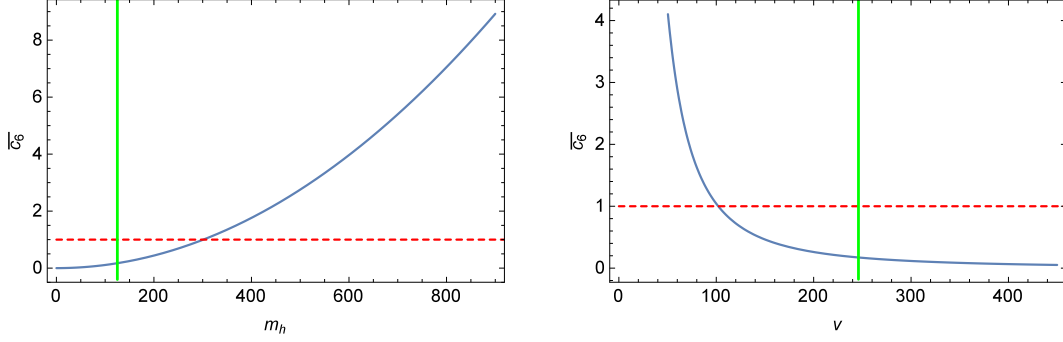
Figure 2. Required value of the Wilson coefficient $\bar{c}_6$ in dependence on the mass of the Higgs boson (left plot) and versus the value of the vev (right plot). See text for details.

thereby eliminating the explicit reference to the new scale Λ (which we shall leave open).

In the left panel of Figure 2 we depict the required value of $\bar{c}_6$ versus the Higgs-boson mass. The red dashed line $\bar{c}_6 \equiv 1$ roughly corresponds to the upper boundary of the viable parameter space, not in conflict with the validity of the perturbative EFT. We can see clearly that a heavy Higgs boson of only $m_h \gtrsim 300\,\text{GeV}$ would have already basically invalidated our approach. The same would have been true for a smaller vev of $v < 100$ GeV (keeping $m_h = 125\,\text{GeV}$), which we show for completeness in the right panel of Figure 2. The experimental values $m_h = 125$ GeV and $v = 246$ GeV, visualized by green vertical lines, are however in perfect agreement with a reasonable value of $\bar{c}_6 \approx 0.17$. Finally, the potential (2.1), employing these values, is plotted as a solid blue line in Figure 3. It exposes the expected mexican-hat form, featuring a stable minimum at a non-trivial field value. We conclude that, while it would have been easily possible that the numerical values of the mass scales generated in nature after EWSB would have excluded our setup, the actual values just lie in a range that allows for EWSB to be triggered by a single $D = 6$ operator instead of a negative mass squared term.

In order for the strict $\mu^2 = 0$ setup to be self-consistent, the μ^2 term should also not be generated again from loop effects within the effective theory, involving physical scales. This is however guaranteed, since the only physical suppression scale Λ can always be factored out of loop integrals and never enters dynamically. Nevertheless, note that it is the UV completion that eventually needs to explain the absence of this term.

Finally, we show that the inclusion of the SM quantum corrections to the potential, generating a term of the form $|H|^4 \log(H^2/\mu_r^2)$ [1], only corresponds to a small perturbation of our setup, leading to little changes in the parameters λ and $\bar{c}_6$. Neglecting the tiny impact of light quarks, the SM contributions to the one-loop
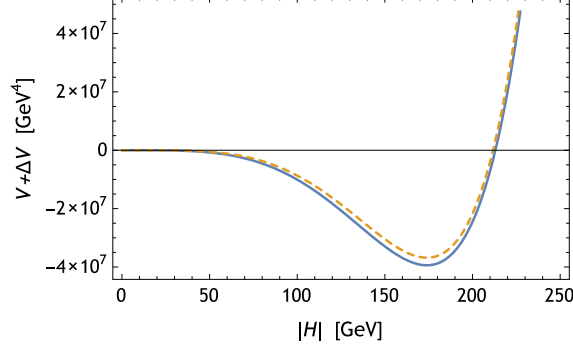
Figure 3. Blue curve: The Higgs potential (2.1), employing the physical values for m_h and v. Red dashed curve: The Higgs potential, including the SM one-loop corrections (3.3), leading to the shifts (3.7).

Coleman-Weinberg potential are given by (see, *e.g.*, [2])

$$\Delta V = \frac{1}{64\pi^2} \sum_{i=W,Z,h,\chi,t} n_i \, M_i^4(H) \left[\log \frac{M_i^2(H)}{\mu_r^2} - C_i \right] . \qquad (3.3)$$

Here, the tree-level field-dependent mass terms read

$$
\begin{aligned}
m_W^2(H) &= \frac{g^2}{2} H^2, \ m_Z^2(H) = \frac{g^2 + g'^2}{2} H^2, \\
m_h^2(H) &= 6\lambda H^2, \ m_\chi^2(H) = 2\lambda H^2, \\
m_t^2(H) &= y_t^2 H^2 ,
\end{aligned}
\qquad (3.4)
$$

where we have dropped contributions suppressed by Λ^2, the numbers of degrees of freedom are

$$n_W = 6, \ n_Z = 3, \ n_h = 1, \ n_\chi = 3, \ n_t = -12, \qquad (3.5)$$

and the constants C_i are given by

$$C_W = C_Z = 5/6, \ C_h = C_\chi = C_t = 3/2 . \qquad (3.6)$$

In the end, the top quark furnishes the dominant correction. Adding (3.3) to (2.1), setting the renormalization scale to $\mu_r = v/\sqrt{2}$, and solving for c_6 and λ that reproduce correctly v and m_h, leads to the shifts

$$\Delta\lambda \approx -0.033, \quad \Delta\bar{c}_6 \approx 0.022 , \qquad (3.7)$$

which is a $\mathcal{O}(10\%)$ effect. We show the resulting potential as a red dashed line in Figure 3. It becomes a little bit flatter before the zero of the undisturbed potential and a bit steeper afterwards. Moreover, there arises a tiny maximum at low values of $|H|$, such that the origin is a minimum - however it corresponds to $(V + \Delta V) = 0$ and thus lies much higher than the global minimum at $|H| = v/\sqrt{2}$. The overall changes are very modest.
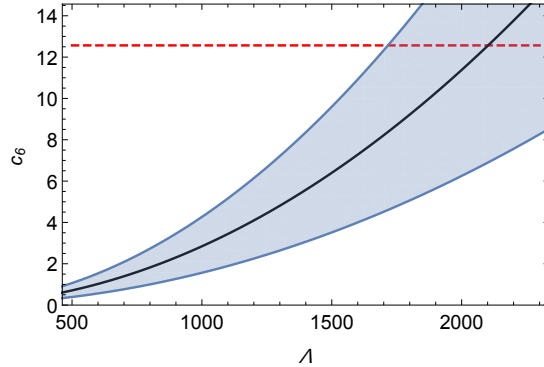
Figure 4. The solid line depicts the required value of the Wilson coefficient c_6 in dependence on the cutoff Λ in our setup, while the blue region allows for an appropriate first order electroweak phase transition such as to trigger electroweak baryogenesis. See text for details.

4 Phenomenology

Beyond the potential direct discovery of new states at (or slightly above) the TeV scale, our model offers distinct signatures in Higgs-pair production and cosmology that we want to discuss in the following.

First of all, the sizable coefficient $\bar{c}_6 \approx 0.2$ leads to a notable change in the production cross section of Higgs pairs, since $\mathcal{O}_6$ contributes to the trilinear Higgs-self interaction after EWSB. In fact, it increases the cross section by a factor of $\sim (2-3)$. This is in a range that should be possible to exclude at the LHC with a luminosity of $\mathcal{L} \gtrsim 600\,\mathrm{fb}^{-1}$, see [3].[3]

Beyond that, the presence of the operator $\mathcal{O}_6$ also modifies the electroweak phase transition. Without this operator, the phase transition is of second order for $m_h = 125$ GeV (see, *e.g.*, [2]). This excludes the possibility of electroweak baryogenesis within the SM as there is no out-of equilibrium dynamics at the phase transition. On the other hand, a sizable contribution of $\mathcal{O}_6$ changes the Higgs potential such that a first order phase transition becomes possible for the physical Higgs mass [4], allowing for electroweak baryogenesis (if enough CP violation is present). In Figure 4 we show again c_6 versus the cutoff Λ, where now the blue region corresponds to a first order phase transition that leads to a stable $T = 0$ minimum, while in addition sphaleron processes are sufficiently suppressed in the broken phase such as to not wash out the generated baryon asymmetry [4]. The latter requirement leads to the condition $\langle h(T_c)\rangle/T_c \gtrsim 1$ at the critical temperature T_c. Very interestingly, our $\mu^2 = 0$ solution just lies in the middle of the preferred region, while the SM (*i.e.*, $c_6 = 0$) does not allow for electroweak baryogenesis.

---

[3]Note that $\bar{c}_6 \approx 0.2$ corresponds to $c_6 \approx -1.2$ in the conventions used to present the final results in [3].

We conclude that the required value of $\bar{c}_6$ leads to a very interesting phenomenology, allowing for pronounced effects in Higgs-pair production as well as opening the possibility of the creation of our current universe via electroweak baryogenesis. This makes the setup avoiding relevant operators attractive on its own. Beyond that, it calls for an examination of how the effective potential (2.1) could be generated - approximately or exactly - from a UV theory. This will be discussed in the Appendix.

5 Conclusions

As we know very little about the dynamics of EWSB or how the hierarchy problem is eventually solved in nature, various approaches to EWSB should be examined and tested, in particular also from the low energy perspective, even if they might not be the most obvious ones. In this article, we have demonstrated that setting the notorious relevant operator $|H|^2$ in the Higgs potential to zero and adding instead an operator $\mathcal{O}_6 = |H|^6$ can lead to a viable electroweak symmetry breaking. We pointed out that it is the lightness of the Higgs-boson that - perhaps unexpectedly - allows for this setup to be self-consistent, in a sense that it is in agreement with perturbativity, allowing a reasonably small c_6 and a natural NP scale of $\Lambda \sim (1 - 2)\,\text{TeV}$. The latter is needed in order to create a mass gap that allows to neglect operators with $D = 8$ and beyond. Moreover, once set to zero, the μ^2 term is not re-generated in the low energy theory. Eliminating this parameter and adding instead the $D = 6$ coefficient c_6 keeps the theory very predictive, since the number of parameters stays the same. In particular, the setup is fully testable in experiments currently under way, since relatively large changes in the Higgs-pair production cross section are predicted.

The main purpose of our analysis was to study the phenomenological viability of this distinct IR limit, allowing new directions in model building. However, in the appendix we will also provide already a first classification of UV completions that could lead to such a setup at least approximately, or exactly, depending on further structures of the UV theory. There, we will also entertain the possibility that, while the Higgs boson could be a fundamental scalar, all scales are generated dynamically.

As it is a distinct theoretical limit, which also opens the possibility of generating our universe via baryogenesis at the electroweak scale and interestingly enough is not ruled out by measurements in the Higgs sector yet, the $\mu^2 \to 0$ model examined here should be considered as an alternative mechanism of breaking electroweak symmetry dynamically.

A   Possible UV Completions

So far, the form of the Higgs potential (2.1) was considered as a matching condition on the unspecified UV completion. Now, we are going to discuss classes of UV models that could generate such a potential at the tree level. The general picture will be that the SM-like theory (featuring $\mu^2 = 0$) possesses no scale at the classical level and is then coupled to a sector that breaks scale invariance, potentially dynamically. Note that such an additional breaking is needed in the first place, since it was found that the breaking of scale invariance within the SM by dimensional transmutation is not sufficient to generate the Higgs mass of $m_h = 125\,\text{GeV}$ (see, *e.g.*, [5, 6]). The NP might itself respect scale invariance at the classical level, generating all masses dynamically. However in the first part of the following discussion we will leave open how the scales in the new sector emerge.

Let us nevertheless already stress a difference of our approach from the usual one, often used in models employing classical scale invariance (CSI) as a building principle. In the latter, the μ^2 term is forbidden at the tree level, but then regenerated spontaneously, usually via the (loop-induced) vev of an additional scalar singlet in a Higgs portal term, mimicking the usual SM Higgs potential (see, *e.g.*, [5, 6], as well as in general [7] on models that generate all scales dynamically). Thus, the low energy phenomenology can be quite similar to the one of the SM. In our approach, however, no relevant operator needs to be generated in the electroweak-scale theory at all. The breaking of scale invariance is induced in an orthogonal - possibly also spontaneous/dynamical - way, via an *irrelevant* operator, introduced by integrating out a heavy field that couples to the SM. This leads to a distinct low energy phenomenology and full testability of our setup. It provides a new minimal way of allowing for a viable EWSB in the presence of the scale-invariant tree-level SM Lagrangian, that interestingly features $m_h \to 0$ in the decoupling limit of $\Lambda \to \infty$. [4]

A.1   General Discussion

We will start with the option of not worrying weather CSI is a good building principle for the full Lagrangian. Depending on how the NP scale is generated, the full models might possess CSI or not.

A scalar field S with mass M_S, transforming as a singlet of $SU(2)_L$, can generate the operator $\mathcal{O}_6$ at the tree level, allowing the NP to be not too light, while a potential contribution to the $|H|^2$ operator could be deferred to the loop level (or beyond, in the presence of additional structures). The singlet can in general feature two interaction structures with the Higgs boson, $M_S \lambda_S S|H|^2$ and $\lambda_p S^2|H|^2$, and both are needed to induce $\mathcal{O}_6$ at the tree level, as shown in Figure 5, leading

---

[4]Also a mixture of both approaches, generating a very small (potentially even positive) μ^2, while assisting with $\mathcal{O}_6$ to trigger EWSB in a theory respecting scale invariance at the tree level, might be interesting to consider.
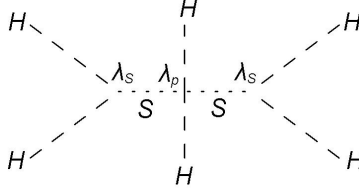
Figure 5. Generation of the operator $\mathcal{O}_6$ by integrating out the new scalar singlet S at the tree level. See text for details.

to $c_6/\Lambda^2 \sim \lambda_p|\lambda_S|^2/M_S^2$. Note in that context that generating c_6 at the loop level would be less straightforward, since unless the perturbativity bound of the underlying theory is saturated, it would in general tend to be too small. If one wants to insist on generating all scales dynamically, the dimensionfull coefficient in front of λ_S could be thought of as a vev of a new field, or arise from a compositeness scale, see below. Since at the tree level only $\mathcal{O}_6$ is generated, one could entertain the possibility that quadratically cutoff-dependent quantum corrections to μ^2 are canceled in an extended singlet NP sector, e.g. by invoking (partial) supersymmetry or a twin Higgs mechanism, such as to approximate, or fully satisfy, (2.1). A related discussion on a complete cancellation of UV effects on the $|H|^2$ operator - which however there is generated again spontaneously - is given in [6] (see also [5]). Alternatively, one could imagine the interaction terms to be cut off by a strong form factor, such that $\lambda_i(p)$ are rapidly vanishing for large momentum p flowing through the vertex. While this would tame loops, it would again not prevent a sizable tree level contribution to $\mathcal{O}_6$ via integrating out S at zero momentum. Finally, one could just enforce the vanishing of the (renormalized) relevant operator by hand at the matching scale. In any case, the scalar S allows to entertain a UV completion where EWSB could be driven by $\mathcal{O}_6$ and not by a negative sign μ^2 term. If the new singlet does not develop a vev and has a mass $M_S \gg v$, with a significant fraction not stemming from the SM-Higgs sector, it should be save from current limits. A survey of the impact of the new fields from the UV completion on Higgs phenomenology (and precision observables) is left for future work. For an overview of constraints on scalar extensions of the SM see, *e.g.*, [8].[5]

A.2 Dynamical Generation of all Scales

Now we are also going to address how the scale in the UV completions emerges, showing concrete possibilities to generate all scales dynamically. It will thus not be necessary to assume a relevant operator at all, neither in the shortest-distance UV theory, nor in the IR limit $E \sim v$. We will first consider compositeness as a rather natural mechanism to obtain a mass from binding energy, without the need of a

---

[5]Extension of the vector-boson sector could also generate $\mathcal{O}_6$ at the tree level, see [9].

relevant operator in the UV theory. Then we will discuss the option of generating a mass from the Coleman-Weinberg mechanism.

A.2.1 Compositeness

One elegant option to obtain the mass for the particles to be integrated out is to consider them to be composites, bound together by a new strong interaction. We assume the condensation scale, beyond which one could resolve the fundamental constituents, to lie well above $1\,\text{TeV}$, such that at currently accessible energies one would never see them directly. In particular, a set of two fermions (or scalars) could form a massive scalar singlet of $SU(2)_L$, given that they transform both in the same representation. These fields could be massless before they condense - however the UV constituents could, depending on their quantum numbers, possibly also get a mass from the Higgs sector after EWSB. Beyond the mass of the composite state, as discussed above, one needs a scale to be present in the $S|H|^2$ interaction. This in turn could also be generated from the composite sector or arise via the vev of a further singlet from the Coleman-Weinberg mechanism (see below).

A.2.2 Coleman-Weinberg Mechanism

Before concluding, let us also entertain the option to generate all scales via the Coleman-Weinberg mechanism. The simplest Coleman-Weinberg setups always rely on a non-negligible vev for a scalar in order to generate its mass [5, 6]. To avoid potential issues with precision tests or Higgs physics from $SU(2)_L$ breaking as well as the direct generation of the μ^2 term, in our case we need to assume the presence of one more scalar particle compared to the general discussion at the beginning of the section, which however could be elementary. In practice, another singlet S_2 could obtain a vev from the Coleman-Weinberg mechanism as explained before (see also [6]). In turn, portal interactions with S would generate its mass dynamically. On the other hand, direct portals to the SM Higgs need to be strongly suppressed in the thought setup. This could *e.g.* be realized geometrically in an extra-dimensional framework, with S being the messengers between the SM and the new singlet S_2, residing on different branes.

Finally, nature might have chosen a completely different way to generate $\mathcal{O}_6$, while avoiding the μ^2 term, still to be found. A further analysis of the potential UV completions, including the examination of dark matter candidates, will be deferred to future work.

References

[1] S. R. Coleman and E. J. Weinberg, Phys. Rev. D **7** (1973) 1888.

[2] M. Quiros, hep-ph/9901312.

[3] F. Goertz, A. Papaefstathiou, L. L. Yang and J. Zurita, arXiv:1410.3471 [hep-ph].

[4] C. Grojean, G. Servant and J. D. Wells, Phys. Rev. D **71** (2005) 036001 [hep-ph/0407019].

[5] R. Hempfling, Phys. Lett. B **379** (1996) 153 [hep-ph/9604278].

[6] C. Englert, J. Jaeckel, V. V. Khoze and M. Spannowsky, JHEP **1304** (2013) 060 [arXiv:1301.4224 [hep-ph]].

[7] W. A. Bardeen, FERMILAB-CONF-95-391-T. J. P. Fatelo, J. M. Gerard, T. Hambye and J. Weyers, Phys. Rev. Lett. **74** (1995) 492. T. Hambye, Phys. Lett. B **371** (1996) 87 [hep-ph/9510266]. P. H. Frampton and C. Vafa, hep-th/9903226. K. A. Meissner and H. Nicolai, Phys. Lett. B **648** (2007) 312 [hep-th/0612165]. W. -F. Chang, J. N. Ng and J. M. S. Wu, Phys. Rev. D **75** (2007) 115016 [hep-ph/0701254]. R. Foot, A. Kobakhidze and R. R. Volkas, Phys. Lett. B **655** (2007) 156 [arXiv:0704.1165]. K. A. Meissner and H. Nicolai, Phys. Lett. B **660** (2008) 260 [arXiv:0710.2840]. S. Iso, N. Okada, Y. Orikasa, Phys. Rev. D **80** (2009) 115007 [arXiv:0909.0128]. M. Shaposhnikov and C. Wetterich, Phys. Lett. B **683** (2010) 196 [arXiv:0912.0208]. L. Alexander-Nunneley and A. Pilaftsis, JHEP **1009** (2010) 021 [arXiv:1006.5916]. T. Hur and P. Ko, Phys. Rev. Lett. **106**, 141802 (2011) [arXiv:1103.2571]. S. Iso and Y. Orikasa, PTEP **2013** (2013) 023B08 [arXiv:1210.2848]. E. J. Chun, S. Jung and H. M. Lee, Phys. Lett. B **725** (2013) 158 [Erratum-ibid. B **730** (2014) 357] [arXiv:1304.5815]. M. Heikinheimo, A. Racioppi, M. Raidal, C. Spethmann and K. Tuominen, Mod. Phys. Lett. A **29** (2014) 1450077 [arXiv:1304.7006]. T. Hambye and A. Strumia, Phys. Rev. D **88** (2013) 055022 [arXiv:1306.2329]. C. D. Carone and R. Ramos, Phys. Rev. D **88** (2013) 055020 [arXiv:1307.8428]. A. Farzinnia, H. J. He and J. Ren, Phys. Lett. B **727** (2013) 141 [arXiv:1308.0295]. E. Gabrielli, M. Heikinheimo, K. Kannike, A. Racioppi, M. Raidal and C. Spethmann, Phys. Rev. D **89** (2014) 1, 015017 [arXiv:1309.6632]. R. Foot, A. Kobakhidze, K. L. McDonald and R. R. Volkas, Phys. Rev. D **89** (2014) 115018 [arXiv:1310.0223]. M. Holthausen, J. Kubo, K. S. Lim and M. Lindner, JHEP **1312** (2013) 076 [arXiv:1310.4423]. C. T. Hill, Phys. Rev. D **89** (2014) 073003 [arXiv:1401.4185]. J. Guo and Z. Kang, arXiv:1401.5609. S. Benic and B. Radovcic, Phys. Lett. B **732** (2014) 91 [arXiv:1401.8183]. P. H. Chankowski, A. Lewandowski, K. A. Meissner and H. Nicolai, Mod. Phys. Lett. A **30** (2015) 02, 1550006 [arXiv:1404.0548]. K. Allison, C. T. Hill and G. G. Ross, Phys. Lett. B **738** (2014) 191 [arXiv:1404.6268]. W. Altmannshofer, W. A. Bardeen, M. Bauer, M. Carena and J. D. Lykken, JHEP **1501** (2015) 032 [arXiv:1408.3429]. O. Antipin, M. Redi and A. Strumia, JHEP **1501** (2015) 157 [arXiv:1410.1817]. G. F. Giudice, G. Isidori, A. Salvio and A. Strumia, JHEP **1502** (2015) 137 [arXiv:1412.2769].

[8] J. de Blas, M. Chala, M. Perez-Victoria and J. Santiago, arXiv:1412.8480 [hep-ph].

[9] F. del Aguila, J. de Blas and M. Perez-Victoria, JHEP **1009** (2010) 033 [arXiv:1005.3998 [hep-ph]].